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# Lepton Flavor Violating Decays, Soft Leptogenesis and SUSY SO(10)

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## Abstract

We investigate lepton flavor violating decays in a SUSY SO(10) model with symmetric textures recently constructed by us. Unlike the models with lop-sided textures which give rise to a large decay rate for $\mu \rightarrow e\gamma$, the decay rate we get is much suppressed and yet it is large enough to be accessible to the next generation of experiments. We have also investigated the possibility of baryogenesis resulting from soft leptogenesis. We find that with the soft SUSY masses assuming their natural values, $B' \equiv \sqrt{BM_1} \sim 1.4\ TeV$ and $Im(A) \sim 1\ TeV$, the observed baryon asymmetry in the Universe can be accommodated in our model. We have also updated the predictions of our model for the masses, mixing angles and CP violating measures in both charged fermion and neutrino sectors, using the most up-to-date experimental data as input.

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## I. INTRODUCTION

After Neutrino 2004, the allowed region for the neutrino oscillation parameters has been reduced significantly, and their measurements have now entered the precision phase. There have been a few supersymmetric (SUSY) $SO(10)$ models constructed aiming to accommodate the observed neutrino masses and mixing angles (For a recent review on SO(10) models, see Ref. [1].) By far, the LMA solution is the most difficult to obtain. Most of the models in the literature assume "lopsided" mass matrices. In our model based on SUSY $SO(10) \times SU(2)$ [2](referred to "CM" herein), we consider *symmetric* mass matrices which result from the left-right symmetric breaking of $SO(10)$ and the breaking of family symmetry $SU(2)$. In view of the much improved experimental data on neutrino oscillation parameters as well as those in the quark mixing from B Physics, we re-analyze our model and find that it can still accommodate all experimental data within $1\sigma$. We investigate several lepton flavor violating (LFV) processes in our model, including the decay of the muon into an electron and a photon, which is the most stringently constrained LFV process. We also investigate in this paper the possibility of baryogenesis utilizing soft leptogenesis.

This paper is organized as follows: In Sec. II, we briefly describe our model, and show its predictions for the masses, mixing angles and CP violating phases in both charged fermion and neutrino sectors, using the most up-to-date experimental data as input. Various decay rates for lepton flavor violation processes are calculated in Sec. III. Sec. IV concerns soft leptogenesis in our model, while Sec. V concludes this paper.

## II. THE MODEL

The details of our model based on $SO(10) \times SU(2)_F$ are contained in CM [2].The following is an outline of its salient features. In order to specify the superpotential uniquely, we invoke $Z_2 \times Z_2 \times Z_2$ discrete symmetry. The matter fields are

$$\psi_a \sim (16, 2)^{-++} \quad (a = 1, 2), \qquad \psi_3 \sim (16, 1)^{+++}$$

where $a = 1, 2$ and the subscripts refer to family indices; the superscripts $+/-$ refer to $(Z_2)^3$ charges. The Higgs fields which break $SO(10)$ and give rise to mass matrices upon acquiring VEV's are

$$\begin{array}{llllll} (10,1): & T_1^{+++}, & T_2^{-+-}, & T_3^{--+}, & T_4^{---}, & T_5^{+--} \\ (\overline{126},1): & & \overline{C}^{---}, & \overline{C}_1^{+++}, & \overline{C}_2^{++-} & \end{array}$$

Higgs representations 10 and $\overline{126}$ give rise to Yukawa couplings to the matter fields which are symmetric under the interchange of family indices. $SO(10)$ is broken through the left-right symmetry breaking chain, and symmetric mass matrices arise. The $SU(2)$ family symmetry [3] is broken in two steps and the mass hierarchy is produced using the Froggatt-Nielsen mechanism: $SU(2) \xrightarrow{\epsilon M} U(1) \xrightarrow{\epsilon' M}$ *nothing* where $M$ is the UV-cutoff of the effective theory above which the family symmetry is exact, and $\epsilon M$ and $\epsilon^{'} M$ are the VEV's accompanying the flavon fields given by

$$\begin{array}{llll} (1,2): & \phi_{(1)}^{++-}, & \phi_{(2)}^{+-+}, & \Phi^{-+-} \\ (1,3): & S_{(1)}^{+--}, & S_{(2)}^{---}, & \Sigma^{++-} \end{array} \tag{1}$$

The various aspects of VEV's of Higgs and flavon fields are given in CM.

The superpotential of our model is

$$W = W_{Dirac} + W_{\nu_{RR}} \tag{2}$$

$$\begin{aligned} W_{Dirac} &= \psi_3\psi_3 T_1 + \frac{1}{M}\psi_3\psi_a\left(T_2\phi_{(1)} + T_3\phi_{(2)}\right) \\ &\quad + \frac{1}{M}\psi_a\psi_b\left(T_4 + \overline{C}\right)S_{(2)} + \frac{1}{M}\psi_a\psi_b T_5 S_{(1)} \\ W_{\nu_{RR}} &= \psi_3\psi_3\overline{C}_1 + \frac{1}{M}\psi_3\psi_a\Phi\overline{C}_2 + \frac{1}{M}\psi_a\psi_b\Sigma\overline{C}_2\ . \end{aligned} \tag{3}$$

The mass matrices then can be read from the superpotential to be

$$\begin{aligned} M_{u,\nu_{LR}} &= \begin{pmatrix} 0 & 0 & \langle 10_2^+\rangle\epsilon' \\ 0 & \langle 10_4^+\rangle\epsilon & \langle 10_3^+\rangle\epsilon \\ \langle 10_2^+\rangle\epsilon' & \langle 10_3^+\rangle\epsilon & \langle 10_1^+\rangle \end{pmatrix} \\ &= \begin{pmatrix} 0 & 0 & r_2\epsilon' \\ 0 & r_4\epsilon & \epsilon \\ r_2\epsilon' & \epsilon & 1 \end{pmatrix} M_U \end{aligned} \tag{4}$$

$$\begin{aligned} M_{d,e} &= \begin{pmatrix} 0 & \langle 10_5^-\rangle\epsilon' & 0 \\ \langle 10_5^-\rangle\epsilon' & (1,-3)\left\langle\overline{126}^-\right\rangle\epsilon & 0 \\ 0 & 0 & \langle 10_1^-\rangle \end{pmatrix} \\ &= \begin{pmatrix} 0 & \epsilon' & 0 \\ \epsilon' & (1,-3)p\epsilon & 0 \\ 0 & 0 & 1 \end{pmatrix} M_D\ , \end{aligned} \tag{5}$$

where $M_U \equiv \langle 10_1^+ \rangle$, $M_D \equiv \langle 10_1^- \rangle$, $r_2 \equiv \langle 10_2^+ \rangle / \langle 10_1^+ \rangle$, $r_4 \equiv \langle 10_4^+ \rangle / \langle 10_1^+ \rangle$ and $p \equiv \left\langle \overline{126}^- \right\rangle / \langle 10_1^- \rangle$. The right-handed neutrino mass matrix is

$$
\begin{aligned}
M_{\nu_{RR}} &= \begin{pmatrix} 0 & 0 & \left\langle \overline{126}_2^{\prime 0} \right\rangle \delta_1 \\ 0 & \left\langle \overline{126}_2^{\prime 0} \right\rangle \delta_2 & \left\langle \overline{126}_2^{\prime 0} \right\rangle \delta_3 \\ \left\langle \overline{126}_2^{\prime 0} \right\rangle \delta_1 & \left\langle \overline{126}_2^{\prime 0} \right\rangle \delta_3 & \left\langle \overline{126}_1^{\prime 0} \right\rangle \end{pmatrix} \\
&= \begin{pmatrix} 0 & 0 & \delta_1 \\ 0 & \delta_2 & \delta_3 \\ \delta_1 & \delta_3 & 1 \end{pmatrix} M_R
\end{aligned}
\tag{6}
$$

with $M_R \equiv \left\langle \overline{126}_1^{\prime 0} \right\rangle$. Here the superscripts $+/-/0$ refer to the sign of the hypercharge. It is to be noted that there is a factor of $-3$ difference between the (22) elements of mass matrices $M_d$ and $M_e$. This is due to the CG coefficients associated with $\overline{126}$; as a consequence, we obtain the phenomenologically viable Georgi-Jarlskog relation. We then parameterize the Yukawa matrices as follows, after removing all the non-physical phases by rephasing various matter fields:

$$
Y_{u,\nu_{LR}} = \begin{pmatrix} 0 & 0 & a \\ 0 & be^{i\theta} & c \\ a & c & 1 \end{pmatrix} d \tag{7}
$$

$$
Y_{d,e} = \begin{pmatrix} 0 & ee^{-i\xi} & 0 \\ ee^{i\xi} & (1,-3)f & 0 \\ 0 & 0 & 1 \end{pmatrix} h \quad . \tag{8}
$$

We use the following as inputs at $M_Z = 91.187\ GeV$ [4, 5]:

$$
\begin{aligned}
m_u &= 2.21\ MeV (2.33^{+0.42}_{-0.45}) \\
m_c &= 682\ MeV (677^{+56}_{-61}) \\
m_t &= 181\ GeV (181^{+}_{-}13) \\
m_e &= 0.486\ MeV (0.486847) \\
m_\mu &= 103\ MeV (102.75) \\
m_\tau &= 1.74\ GeV (1.7467) \\
|V_{us}| &= 0.225 (0.221 - 0.227) \\
|V_{ub}| &= 0.00368 (0.0029 - 0.0045) \\
|V_{cb}| &= 0.0392 (0.039 - 0.044)
\end{aligned}
$$

where the values extrapolated from experimental data are given inside the parentheses. Note that the masses given above are defined in the modified minimal subtraction ($\overline{\mathrm{MS}}$) scheme and are evaluated at $M_Z$. These values correspond to the following set of input parameters at the GUT scale, $M_{GUT} = 1.03 \times 10^{16}\ GeV$:

$$
\begin{gathered}
a = 0.00250, \quad b = 3.26 \times 10^{-3} \\
c = 0.0346, \quad d = 0.650 \\
\theta = 0.74 \\
e = 4.036 \times 10^{-3}, \quad f = 0.0195 \\
h = 0.06878, \quad \xi = -1.52 \\
g_1 = g_2 = g_3 = 0.746
\end{gathered} \tag{9}
$$

the one-loop renormalization group equations for the MSSM spectrum with three right-handed neutrinos are solved numerically down to the effective right-handed neutrino mass scale, $M_R$. At $M_R$, the seesaw mechanism is implemented. With the constraints $|m_{\nu_3}| \gg |m_{\nu_2}|, |m_{\nu_1}|$ and maximal mixing in the atmospheric sector, the up-type mass texture leads us to choose the following effective neutrino mass matrix

$$
M_{\nu_{LL}} = \begin{pmatrix} 0 & 0 & t \\ 0 & 1 & 1+t^n \\ t & 1+t^n & 1 \end{pmatrix} \frac{d^2 v_u^2}{M_R} \tag{10}
$$

with $n = 1.15$, and from the seesaw formula we obtain

$$
\delta_1 = \frac{a^2}{r} \tag{11}
$$

$$
\delta_2 = \frac{b^2 t e^{2i\theta}}{r} \tag{12}
$$

$$
\delta_3 = \frac{-a(be^{i\theta}(1+t^{1.15}) - c) + bcte^{i\theta}}{r}, \tag{13}
$$

where $r = (c^2 t + a^2 t^{0.15}(2 + t^{1.15}) - 2a(-1 + c + ct^{1.15}))$. We then solve the two-loop RGE's for the MSSM spectrum down to the SUSY breaking scale, taken to be $m_t(m_t) = 176.4\ GeV$, and then the SM RGE's from $m_t(m_t)$ to the weak scale, $M_Z$. We assume that $\tan\beta \equiv v_u/v_d = 10$, with $v_u^2 + v_d^2 = (246/\sqrt{2}\ GeV)^2$. At the weak scale $M_Z$, the predictions for $\alpha_i \equiv g_i^2/4\pi$ are

$$
\alpha_1 = 0.01663, \quad \alpha_2 = 0.03374, \quad \alpha_3 = 0.1242\ .
$$

These values compare very well with the values extrapolated to $M_Z$ from the experimental data, $(\alpha_1, \alpha_2, \alpha_3) = (0.01696, 0.03371, 0.1214 \pm 0.0031)$. The predictions at the weak scale $M_Z$ for the

TABLE I: The predictions for the charged fermion masses, the CKM matrix elements and the CP violation measures.

| | experimental results extrapolated to $M_Z$ | predictions at $M_z$ |
|---|---|---|
| $m_s/m_d$ | $17 \sim 25$ | 25 |
| $m_s$ | $93.4^{+11.8}_{-13.0} MeV$ | $86.0 MeV$ |
| $m_b$ | $3.00 \pm 0.11 GeV$ | $3.03 GeV$ |
| $\|V_{ud}\|$ | $0.9739 - 0.9751$ | 0.974 |
| $\|V_{cd}\|$ | $0.221 - 0.227$ | 0.225 |
| $\|V_{cs}\|$ | $0.9730 - 0.9744$ | 0.973 |
| $\|V_{td}\|$ | $0.0048 - 0.014$ | 0.00801 |
| $\|V_{ts}\|$ | $0.037 - 0.043$ | 0.0386 |
| $\|V_{tb}\|$ | $0.9990 - 0.9992$ | 0.999 |
| $J^q_{CP}$ | $(2.88 \pm 0.33) \times 10^{-5}$ | $2.87 \times 10^{-5}$ |
| $\sin 2\alpha$ | $-0.16 \pm 0.26$ | $-0.048$ |
| $\sin 2\beta$ | $0.736 \pm 0.049$ | 0.740 |
| $\gamma$ | $60^0 \pm 14^0$ | $64^0$ |
| $\overline{\rho}$ | $0.20 \pm 0.09$ | 0.173 |
| $\overline{\eta}$ | $0.33 \pm 0.05$ | 0.366 |

charged fermion masses, CKM matrix elements and strengths of CP violation, are summarized in Table. I. The predictions of our model in this *updated* fit are in good agreement with all experimental data within $1\sigma$, including much improved measurements in B Physics that give rise to precise values for the CKM matrix elements and for the unitarity triangle [6]. Note that we have taken the SUSY threshold correction to $m_b$ to be $-18\%$ [7].

The allowed region for the neutrino oscillation parameters has been reduced significantly after Neutrino 2004. In the atmospheric sector, the global analysis including the most recent K2K result yields, at 90% CL [8],

$$\Delta m^2_{atm} = 2.3^{+0.7}_{-0.4} \times 10^{-3} eV^2 \tag{14}$$

$$\sin^2 2\theta_{atm} > 0.9 \tag{15}$$

$$(\text{best fit value:} \sin^2 2\theta_{atm} = 1.0) \quad . \tag{16}$$

In the solar sector, the global analysis with SNO and most recent KamLAND data yields, at

$1\sigma$ $(3\sigma)$ [9],

$$\Delta m^2_\odot = 8.2^{+0.3}_{-0.3}(^{+1.0}_{-0.8}) \times 10^{-5} eV^2 \tag{17}$$

$$\tan^2\theta_\odot = 0.39^{+0.05}_{-0.04}(^{+0.19}_{-0.11}) \quad . \tag{18}$$

Combining with the CHOOZ result, a global analysis shows that the angle $\theta_{13}$ is constrained to be [9]

$$\sin^2\theta_{13} < 0.015(0.048) \tag{19}$$

at $1\sigma$ $(3\sigma)$. Using the mass square difference in the atmospheric sector $\Delta m^2_{atm} = 2.33 \times 10^{-3}\ eV^2$ and the mass square difference for the LMA solution $\Delta m^2_\odot = 8.14 \times 10^{-5}\ eV^2$ as input parameters, we determine $t = 0.344$ and $M_R = 6.97 \times 10^{12} GeV$, which yield $(\delta_1, \delta_2, \delta_3) = (0.00120, 0.000703e^{i\ (1.47)}, 0.0210e^{i\ (0.175)})$. We obtain the following predictions in the neutrino sector: The three mass eigenvalues are give by

$$(m_{\nu_1}, m_{\nu_2}, m_{\nu_3}) = (0.00262, 0.00939, 0.0492)\ eV\ . \tag{20}$$

The prediction for the MNS matrix is

$$|U_{MNS}| = \begin{pmatrix} 0.852 & 0.511 & 0.116 \\ 0.427 & 0.560 & 0.710 \\ 0.304 & 0.652 & 0.695 \end{pmatrix} \tag{21}$$

which translates into the mixing angles in the atmospheric, solar and reactor sectors,

$$\sin^2 2\theta_{atm} \equiv \frac{4|U_{\mu\nu_3}|^2|U_{\tau\nu_3}|^2}{(1-|U_{e\nu_3}|^2)^2} = 1.00 \tag{22}$$

$$\tan^2\theta_\odot \equiv \frac{|U_{e\nu_2}|^2}{|U_{e\nu_1}|^2} = 0.36 \tag{23}$$

$$\sin^2\theta_{13} = |U_{e\nu_3}|^2 = 0.0134\ . \tag{24}$$

The prediction of our model for the strengths of CP violation in the lepton sector are

$$J^l_{CP} \equiv Im\{U_{11}U^*_{12}U^*_{21}U_{22}\} = -0.00941 \tag{25}$$

$$(\alpha_{31}, \alpha_{21}) = (0.934, -1.49)\ . \tag{26}$$

Using the predictions for the neutrino masses, mixing angles and the two Majorana phases, $\alpha_{31}$ and $\alpha_{21}$, the matrix element for the neutrinoless double $\beta$ decay can be calculated and is given by

$| < m > | = 3.1 \times 10^{-3}\ eV$, with the present experimental upper bound being $0.35\ eV$ [4]. Masses of the heavy right-handed neutrinos are

$$M_1 \;=\; 1.09 \times 10^{7}\ GeV \tag{27}$$

$$M_2 \;=\; 4.53 \times 10^{9}\ GeV \tag{28}$$

$$M_3 \;=\; 6.97 \times 10^{12}\ GeV\ . \tag{29}$$

The prediction for the $\sin^2\theta_{13}$ value is 0.0134, in agreement with the current bound 0.015 at $1\sigma$. Because our prediction for $\sin^2\theta_{13}$ is very close to the present sensitivity of the experiment, the validity of our model can be tested in the foreseeable future [10].

## III. LEPTON FLAVOR VIOLATING DECAYS

In light of the neutrino oscillation, extensive searches for lepton flavor violation processes, such as $\ell_i \to \ell_j\gamma$, $\ell_i^- \to \ell_j^-\ell_j^+\ell_j^-$, muon-electron conversion, are underway. In the SM, as the lepton number is conserved, there is no lepton flavor violation. Non-zero neutrino masses imply lepton number violation. If neutrino masses are induced by the seesaw mechanism, new Yukawa coupling involving the RH neutrinos can induce flavor violation [11], similar to its quark counter part. In the non-supersymmetric case, the decay amplitudes for these processes are inversely proportional to the RH neutrino mass, $M_R^2$, which is typically much higher than the electroweak scale. As a consequence, in non-supersymmetric models, these processes are highly suppressed to the level that are unobservable.

Significant enhancement in the decay rate can be obtained in supersymmetric models, as the characteristic scale in this case is the SUSY scale, which is expected to be not too far from the electroweak scale. Thus the amplitudes for these decay processes scale as inverse square of the SUSY breaking scale, rather than $1/M_R^2$. The relevant interactions that give rise to lepton flavor violating decays come from the soft-SUSY breaking Lagrangian,

$$\begin{aligned} -\mathcal{L}_{soft} \;=\; & (m_{\widetilde{L}}^2)_{ij}\widetilde{\ell}_{L_i}^\dagger\widetilde{\ell}_{L_j} + (m_{\widetilde{e}}^2)_{ij}\widetilde{e}_{R_i}^\dagger\widetilde{e}_{R_j} + (m_{\widetilde{\nu}}^2)_{ij}\widetilde{\ell}_{R_i}^\dagger\widetilde{\ell}_{R_j} \\ & +(\widetilde{m}_{h_d}^2)\widetilde{H}_d^\dagger\widetilde{H}_d + (\widetilde{m}_{h_2}^2)\widetilde{H}_u^\dagger\widetilde{H}_u + \Big[A_\nu^{ij}\widetilde{H}_u\widetilde{\nu}_{R_i}^*\widetilde{\nu}_{L_j} \\ & +A_e^{ij}H_d\widetilde{e}_{R_i}^*\widetilde{e}_{L_j} + \frac{1}{2}B_\nu^{ij}\widetilde{\nu}_{R_i}\widetilde{\nu}_{R_j} + B_hH_dH_u \\ & +h.c.\Big]\quad , \end{aligned} \tag{30}$$

where $\widetilde{\ell}_L$, $\widetilde{e}_R$ and $\widetilde{\nu}_R$ are the LH slepton doublets, RH charged sleptons, and RH sneutrinos, respectively; $H_u$ ($\widetilde{H}_u$) and $H_d$ ($\widetilde{H}_d$) are the two Higgs (higgsino) doublets in MSSM. Assuming

mSUGRA boundary conditions at the GUT scale,

$$(m^2_{\widetilde{L}})_{ij} = (m^2_{\widetilde{e}})_{ij} = (m^2_{\widetilde{\nu}})_{ij} = m_0 \delta_{ij} \tag{31}$$

$$\widetilde{m}^2_{H_d} = \widetilde{m}^2_{H_u} = m_0^2 \tag{32}$$

$$A^{ij}_\nu = (Y_\nu)_{ij} A_0, \quad A^{ij}_e = (Y_e)_{ij} A_0 \tag{33}$$

$$B^{ij}_\nu = M_{\nu_{RR}} B_0, \quad B_h = \mu B_0 \tag{34}$$

where $Y_\nu$ and $Y_e$ are the Yukawa couplings of the neutrinos and charged leptons, and $M_{\nu_{RR}}$ is the Majorana mass matrix of the RH neutrinos. As the slepton mass matrix $(m^2_{\widetilde{L}})_{ij}$ is flavor-blind at the GUT scale, there is no flavor violation at $M_{GUT}$. However, as $(m^2_{\widetilde{L}})_{ij}$ evolves from $M_{GUT}$ to the RH neutrino mass scale, $M_R$, according to the renormalization group equation,

$$\begin{aligned} \frac{d}{d \ln \mu} (m^2_{\widetilde{L}})_{ij} &= \frac{1}{16\pi^2} \Big[ m^2_{\widetilde{L}} (Y^\dagger_\nu Y_\nu)_{ij} \\ &\quad + 2 \Big( (Y^\dagger_\nu m^2_{\widetilde{\nu}} Y_\nu)_{ij} + m^2_{\widetilde{h}} (Y^\dagger_\nu Y_\nu)_{ij} \\ &\quad + (A^\dagger_\nu A_\nu)_{ij} \Big) \Big] , \quad \text{for } i \neq j , \end{aligned} \tag{35}$$

the off diagonal elements in the slepton mass matrix $m^2_{\widetilde{L}}$ can be generated at low energies due to the RG corrections [12],

$$\begin{aligned} \delta(m^2_{\widetilde{L}})_{ij} &= -\frac{1}{8\pi} (3m_0^2 + A_0^2) \\ &\quad \times \sum_{k=1,2,3} (\mathcal{Y}^\dagger_\nu)_{ik} (\mathcal{Y}_\nu)_{kj} \ln(\frac{M_{GUT}}{M_{R_k}}) , \end{aligned} \tag{36}$$

for $i \neq j$. Here $\mathcal{Y}_\nu$ is the Yukawa couplings for the neutrinos in the basis where both charged lepton Yukawa matrix and the Majorana mass matrix for the RH neutrinos are diagonal; $M_{R_k}$ are the masses of the heavy neutrinos. The Yukawa coupling $\mathcal{Y}_\nu$ in the new basis is related to $Y_\nu$ in the original basis by

$$\mathcal{Y}_\nu = P_R O_R Y_\nu O^\dagger_{e_L} . \tag{37}$$

Here $O_{L_e}$ is the diagonalization matrix for

$$\mathcal{M}^{\mathrm{diag}}_e = O_{e_R} M_e O^\dagger_{e_L} , \tag{38}$$

and the diagonal phase matrix $P_R$ and the orthogonal matrix $O_R$ are defined by,

$$\begin{aligned} \mathcal{M}^{\mathrm{diag}}_{\nu_{RR}} &= \mathrm{diag}(M_1, M_2, M_3) \\ &= P_R O_R M_{\nu_{RR}} O^T_R P_R , \end{aligned} \tag{39}$$

TABLE II: Summary of current status and future proposals of the experimental searches for lepton flavor violating decays.

| Decay | current bound on the branching ratio | reach of future experiment |
|---|---|---|
| $\mu \to e\gamma$ | $< 1.2 \times 10^{-11}$ (MEGA, 1999)[13] | $10^{-14}$ (PSI)[14] |
| | | $10^{-15}$ (J-PARC) |
| $\mu \to 3e$ | $< 1.0 \times 10^{-12}$ (SINDRUM, 1988)[15] | |
| $\mu \to e$ in ${}^{48}_{22}Ti$ | $< 6.1 \times 10^{-13}$ (SINDRUM II, 1998)[16] | $2.0 \times 10^{-17}$ (MECO)[17] |
| | | $10^{-18}$ (J-PARC) |
| $\tau \to \mu\gamma$ | $< 3.1 \times 10^{-7}$ (BELLE, 2003) [18] | $10^{-9}$ (BELLE)[18] |
| $\tau \to e\gamma$ | $< 3.6 \times 10^{-7}$ (BELLE, 2003) [19] | |

where $M_{1,2,3}$ are real and positive, and their numerical values are given in Eq. (27)-(29). In our model, the Yukawa matrix $\mathcal{Y}_\nu$ is,

$$\mathcal{Y}_\nu = \begin{pmatrix} 2.69 \times 10^{-6} e^{-(0.695)i} & 5.92 \times 10^{-5} e^{-(2.75)i} & 6.54 \times 10^{-4} e^{-(1.68)i} \\ 1.44 \times 10^{-4} e^{(1.54)i} & 1.73 \times 10^{-3} e^{-(0.176)i} & 8.91 \times 10^{-3} e^{-(1.32)i} \\ 2.18 \times 10^{-3} e^{(0.737)i} & 0.0213 e^{(0.0064)i} & 0.618 \end{pmatrix} . \quad (40)$$

The non-vanishing off-diagonal matrix elements in $(\delta m^2_{\widetilde{L}})_{ij}$ induces lepton flavor violating processes mediated by the superpartners of the neutrinos through the one-loop diagram shown in Fig. 1.

In Table II we summarize current status and future proposals of the experimental searches for lepton flavor violating decays. In the following subsections, we discuss each LFV process individually.

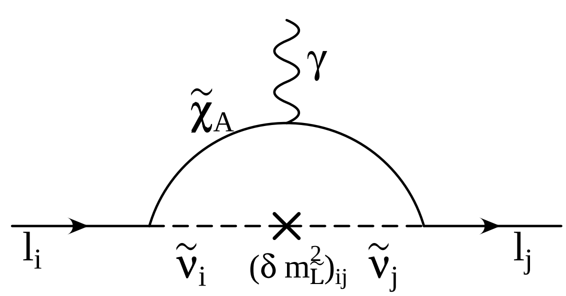


FIG. 1: The dominant diagram that contribute to the decay $\ell_i \to \ell_j\gamma$ at one loop, mediated by the neutralino $\widetilde{\chi}_A$ and the sneutrinos $\widetilde{\nu}$. The inserted mass term $(\delta m^2_{\widetilde{L}})_{ij}$ is induced by the renormalization group evolution from the GUT scale to the RH neutrino mass scales.

### A. $\mu \to e\gamma$, $\tau \to \mu\gamma$ and $\tau \to e\gamma$

The branching ratios for the decay of $\ell_i \to \ell_j + \gamma$ induced by the renormalization group effects described above is given by [12]

$$\begin{aligned} Br(\ell_i \to \ell_j \gamma) &= \frac{\alpha^3}{G_F^2 m_S^8} \left| \frac{-1}{8\pi}(3m_0^2 + A_0^2) \right|^2 \tan^2\beta \\ &\times \left| \sum_{k=1,2,3} (\mathcal{Y}_\nu^\dagger)_{ik} (\mathcal{Y}_\nu)_{kj} \ln\left(\frac{M_{GUT}}{M_{R_k}}\right) \right|^2 \quad . \end{aligned} \tag{41}$$

Here $\alpha$ is the fine structure constant, $G_F$ is the Fermi constant, and $m_S$ is the typical SUSY scalar mass which is given by, to a very good approximation [20],

$$m_S^8 = \frac{1}{2} m_0^2 M_{1/2}^2 (m_0^2 + 0.6 M_{1/2}^2)^2, \tag{42}$$

where $M_{1/2}$ is the universal gaugino mass. In our model, $|\delta(m_{\tilde{L}}^2)_{ij}|$ is given by,

$$\begin{aligned} |\delta(m_{\tilde{L}}^2)_{ij}| &= \left| \frac{1}{8\pi}(3m_0^2 + A_0^2) \right| \\ &\times \begin{pmatrix} * & 3.41 \times 10^{-4} & 0.0098 \\ 3.41 \times 10^{-4} & * & 0.0962 \\ 0.0098 & 0.0962 & * \end{pmatrix} , \end{aligned} \tag{43}$$

for $i \neq j$. Thus the following relation is predicted,

$$Br(\mu \to e\gamma) < Br(\tau \to e\gamma) < Br(\tau \to \mu\gamma) \ . \tag{44}$$

Similar relation was observed in Ref. [21] in which symmetric mass matrices with four texture zeros are utilized. We also note that the value for $\tan\beta$ is 10, thus there is no $\tan\beta$ enhancement in our predictions.

Currently the most stringent experimental bound on the lepton flavor violating processes is on the decay $\mu \to e\gamma$. The prediction of our model for $Br(\mu \to e\gamma)$ is well below the most stringent bound up-to-date from MEGA at LANL [13]. In Fig. 2, the branching ratio of the decay $\mu \to e\gamma$ as a function of the universal gaugino mass $M_{1/2}$ is shown for various scalar masses $A_0$ and $m_0$. For large $A_0$ and low $m_0$ and $M_{1/2}$, there is a large soft SUSY parameter space that give rise to predictions which can be probed by MEG at PSI and/or at J-PARC. In Fig. 3, the branching ratio of the decay $\tau \to \mu\gamma$ as a function of the universal gaugino mass $M_{1/2}$ is shown for various

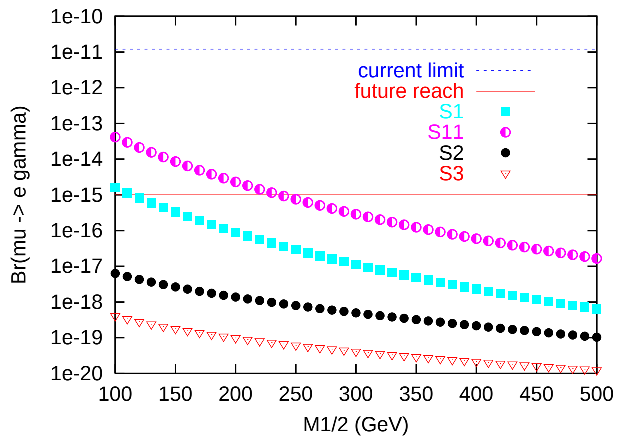


FIG. 2: The branching ratio of the decay $\mu \to e\gamma$ as a function of the universal gaugino mass $M_{1/2}$ for various scalar masses $A_0$ and $m_0$. (S1): $m_0 = A_0 = 100\ GeV$; (S11): $m_0 = 100\ GeV, A_0 = 1\ TeV$; (S2): $m_0 = A_0 = 500\ GeV$; (S3): $m_0 = A_0 = 1\ TeV$. The dash line corresponds to the current experimental limit $1.2 \times 10^{-11}$ from MEGA, while the solid line indicates the reach of a future experiment at J-PARC, $10^{-15}$. The value of $\tan\beta$ in our model is $\tan\beta = 10$.

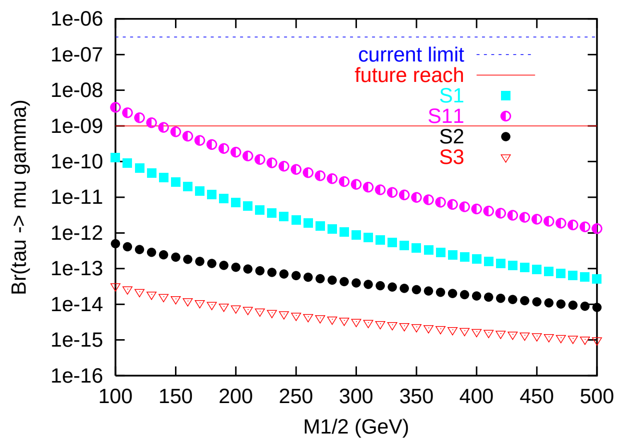


FIG. 3: The branching ratio of the decay $\tau \to \mu\gamma$ as a function of the universal gaugino mass $M_{1/2}$ for various scalar masses $A_0$ and $m_0$. (S1): $m_0 = A_0 = 100\ GeV$; (S11): $m_0 = 100\ GeV, A_0 = 1\ TeV$; (S2): $m_0 = A_0 = 500\ GeV$; (S3): $m_0 = A_0 = 1\ TeV$. The dash line corresponds to the current experimental limit $3.1 \times 10^{-7}$ from BELLE, while the solid line indicates the reach of a future experiment at BELLE, $10^{-9}$. The value of $\tan\beta$ in our model is $\tan\beta = 10$.

scalar masses $A_0$ and $m_0$. For $A_0 \sim \mathcal{O}(1\ TeV)$ and $m_0$ and $M_{1/2}$ both of order $\mathcal{O}(100\ GeV)$, the prediction of our model on $\tau \to \mu\gamma$ may be tested at BELLE in the future. In Fig. 4, the branching ratio of the decay $\tau \to e\gamma$ as a function of the universal gaugino mass $M_{1/2}$ is shown for various scalar mass $A_0$ and $m_0$. For the SUSY parameter space we consider, the prediction for $Br(\tau \to e\gamma)$ is at least four orders of magnitudes below the current experimental upper bound.

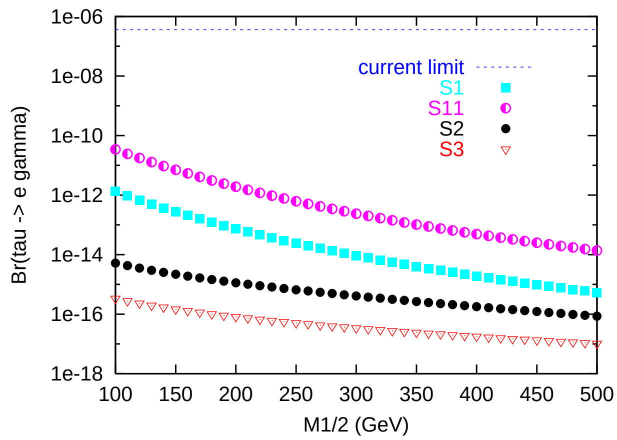


FIG. 4: The branching ratio of the decay $\tau \to e\gamma$ as a function of the universal gaugino mass $M_{1/2}$ for various scalar masses $A_0$ and $m_0$. (S1): $m_0 = A_0 = 100\ GeV$; (S11): $m_0 = 100\ GeV, A_0 = 1\ TeV$; (S2): $m_0 = A_0 = 500\ GeV$; (S3): $m_0 = A_0 = 1\ TeV$. The dash line corresponds to the current upper bound, $3.6 \times 10^{-7}$, from BELLE. The value of $\tan\beta$ in our model is $\tan\beta = 10$.

We comment that, in models with lop-sided textures [22], the maximal mixing angle observed in the atmospheric neutrino sector is due to a large (23) mixing in the charged lepton sector. As a result, the off-diagonal elements in (23) sector of $O_{e_L}$ are of order $\mathcal{O}(1)$, which in turn gives rise to an enhancement in the decay branching ratios. In order to satisfy the current experimental upper bound, some new mechanism must be in place to suppress the decay rate of $\mu \to e\gamma$ in models with lop-sided textures [23]. In our model which utilizes symmetric textures, as large leptonic mixing in our model is a result of the seesaw mechanism, all off-diagonal matrix elements in $Y_\nu$, $O_{e_L}$ and $O_R$ are much smaller than unity, leading to a much smaller branching ratio for $\mu \to e\gamma$ than that predicted in models with lop-sided textures. Yet our prediction is large enough to be probed by the next generation of experiments within a few years.

### B. $\mu \to 3e$

For the process $\mu \to 3e$, as penguium diagrams are the dominant contributions, the branching ratio of the decay $\ell_i^- \to \ell_j^- \ell_j^+ \ell_j^-$ has similar structure as that of the decay $\ell_i^- \to \ell_j^- \gamma$. To a very good approximation, the relation between these two processes reads [12],

$$\frac{Br(\ell_i^- \to \ell_j^- \ell_j^+ \ell_j^-)}{Br(\ell_i^- \to \ell_j^- \gamma)} \simeq \frac{\alpha}{8\pi}\left[\frac{16}{3}\ln(\frac{m_{\ell_i}}{2m_{\ell_j}}) - \frac{14}{9}\right] \quad , \tag{45}$$

where $m_{\ell_i}$ is the $i-$th generation lepton mass. For the decay $\mu \to 3e$, we thus have

$$Br(\mu \to 3e) \simeq 7 \times 10^{-3} Br(\mu \to e\gamma)\ . \tag{46}$$

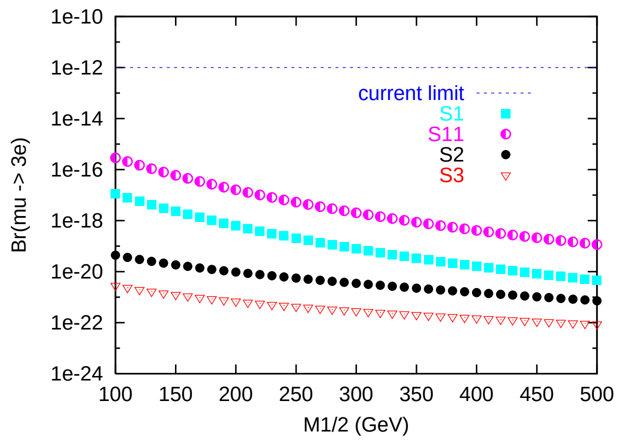


FIG. 5: The branching ratio of the decay $\mu^- \to e^- e^+ e^-$ as a function of the universal gaugino mass $M_{1/2}$ for various scalar masses $A_0$ and $m_0$. (S1): $m_0 = A_0 = 100\ GeV$; (S11): $m_0 = 100\ GeV, A_0 = 1\ TeV$; (S2): $m_0 = A_0 = 500\ GeV$; (S3): $m_0 = A_0 = 1\ TeV$. The dash line corresponds to the current experimental limit $1.0 \times 10^{-12}$ from SINDRUM. The value of $\tan\beta$ in our model is $\tan\beta = 10$.

In Fig. 5, the branching ratio of the decay $\mu \to 3e$ as a function of the universal gaugino mass $M_{1/2}$ is shown for various scalar mass $A_0$ and $m_0$. As the current experimental upper bound and the reach of the next phase of experiment at BELLE are still quite high, the prediction for $\mu \to 3e$ in our model can not be tested, even with a high value of the scalar mass, $A_0 = 1\ TeV$.

### C. $\mu$-$e$ Conversion

Similar to the case of $\mu \to 3e$, the branching ratio for muon-electron conversion is also related to the branching ratio of the decay $\mu \to e\gamma$ as long as $\tan\beta$ is not too small. In the region $\tan\beta > 1$, the relation between these two processes is given by, to a very good approximation [12],

$$\frac{Br(\mu \to e)}{Br(\mu \to e\gamma)} \simeq 16\alpha^4 Z_{eff}^4 Z |F(q^2)|^2 \quad , \tag{47}$$

where $Z_{eff}$ is the effective charge of the nucleon, $Z$ is the proton number and $F(q^2)$ is the nuclear form factor at momentum transfer $q$. For ${}^{48}_{22}Ti$, the conversion rate is

$$Br(\mu \to e;\ {}^{48}_{22}Ti) \simeq 6 \times 10^{-3} Br(\mu \to e\gamma) \quad , \tag{48}$$

where $Z_{eff} = 17.6$ and $F(q^2 = -m_\mu^2) = 0.54$ have been used. In Fig. 6, the branching ratio of the decay $\mu \to e$ in ${}^{48}_{22}Ti$ as a function of the universal gaugino mass $M_{1/2}$ is shown for various scalar mass $A_0$ and $m_0$. For low values of $m_0$ and $M_{1/2}$, there is a very large soft SUSY parameter space that give rise to prediction for $\mu - e$ conversion rate that is sensitive to MECO [17] at BNL and the proposal at J-PARC.

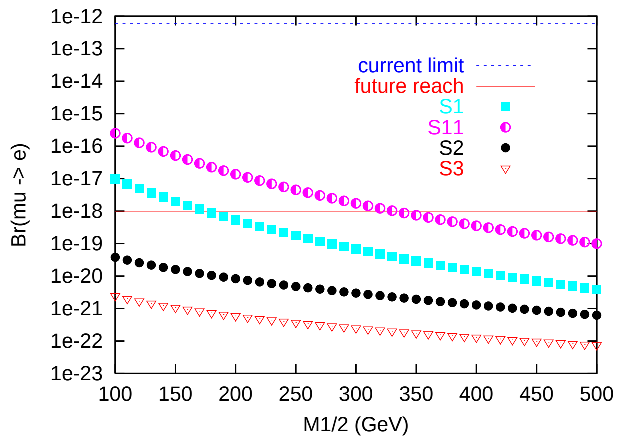


FIG. 6: The branching ratio of the decay $\mu^- \to e^-$ in ${}^{48}_{22}$Ti. as a function of the universal gaugino mass $M_{1/2}$ for various scalar masses $A_0$ and $m_0$. (S1): $m_0 = A_0 = 100\ GeV$; (S11): $m_0 = 100\ GeV, A_0 = 1\ TeV$; (S2): $m_0 = A_0 = 500\ GeV$; (S3): $m_0 = A_0 = 1\ TeV$. The dash line corresponds to the current experimental limit $6.1 \times 10^{-13}$ from SINDRUM II, while the solid line indicates the reach of a future experiment at J-PARC, $10^{-18}$. The value of $\tan\beta$ in our model is $\tan\beta = 10$.

## IV. BARYOGENESIS À *LA* SOFT LEPTOGENESIS

It is well known that the CP violation in the quark sector is too small to explain the observed baryon asymmetry of the Universe (BAU), expressed in terms of the ratio of the baryon number to entropy [24],

$$\frac{n_b}{s} = (0.87 \pm 0.04) \times 10^{-10} \,, \tag{49}$$

derived from CMB and nucleosynthesis measurements. In leptogenesis, leptonic CP violating phases are used to produce asymmetry in leptonic number which then is converted into baryon asymmetry by the electroweak non-perturbative effects due to sphalerons. There are two ways of producing lepton number asymmetry: (i) Standard leptogenesis (STDL) [25] and (ii) Soft leptogenesis (SFTL) [26, 27, 28].

In STDL scenario, the primordial leptonic asymmetry is generated by the decay of the heavy right-handed Majorana neutrinos and their scalar partners, mediated by the Yukawa interactions in the superpotential. In our model, the large hierarchy among the three heavy neutrinos leads to a very small CP asymmetry, which is of the order of $\mathcal{O}(10^{-9})$. In addition, the low value for the mass of the lightest RH neutrino, $M_1 = 1.09 \times 10^7\ GeV$, leads to an extremely large wash-out effect. Due to these reasons, the prediction in our model for the baryonic asymmetry utilizing the standard leptogenesis is of the order of $\mathcal{O}(10^{-15})$, which is four orders of magnitude below the value derived from experimental observations.

SFTL utilizes the soft SUSY breaking sector, and the asymmetry in the lepton number is generated in the decay of the superpartner of the RH neutrinos [26, 27], as opposed to the lightest RH neutrino in the case of STDL. Unlike in STDL where the Yukawa sector is responsible for the required CP violation and lepton number violation, in the scenario of SFTL, the CP violation and lepton number violation trace their origins to SUSY breaking. As a result, it allows a much lower bound on the mass of the lightest RH neutrino, $M_1$, compared to that in STDL. In fact, it has been shown very recently that in contrast to the STDL scenario in which $M_1 > 10^9 GeV$ is typically required to have sufficient baryonic asymmetry [29], SFTL can only work in the region where $M_1 < 10^9\ GeV$ [30]. As a result, the problem of the gravitino over-production [31] may be avoided.

For SFTL, the relevant soft SUSY Lagrangian that involves lightest RH sneutrinos $\widetilde{\nu}_{R_1}$ is the following,

$$\begin{aligned} -\mathcal{L}_{soft} &= (\frac{1}{2}BM_1\widetilde{\nu}_{R_1}\widetilde{\nu}_{R_1} + A\mathcal{Y}_{1i}\widetilde{L}_i\widetilde{\nu}_{R_1}H_u + h.c.) \\ &\quad +\widetilde{m}^2\widetilde{\nu}^\dagger_{R_1}\widetilde{\nu}_{R_1} \ . \end{aligned} \tag{50}$$

This soft SUSY Lagrangian and the superpotential that involves the lightest RH neutrino, $N_1$,

$$W = M_1 N_1 N_1 + \mathcal{Y}_{1i} L_i N_1 H_u \tag{51}$$

give rise to the following interactions

$$\begin{aligned} -\mathcal{L}_{\mathscr{A}} &= \widetilde{\nu}_{R_1}(M_1 Y^*_{1i}\widetilde{\ell}^*_i H^*_u + \mathcal{Y}_{1i}\overline{\widetilde{H}}_u \ell^i_L + A\mathcal{Y}_{1i}\widetilde{\ell}_i H_u) \\ &\quad +h.c. \quad , \end{aligned} \tag{52}$$

and mass terms (to leading order in soft SUSY breaking terms),

$$-\mathcal{L}_{\mathscr{M}} = (M_1^2\widetilde{\nu}^\dagger_{R_1}\widetilde{\nu}_{R_1} + \frac{1}{2}BM_1\widetilde{\nu}_{R_1}\widetilde{\nu}_{R_1}) + h.c. \ . \tag{53}$$

Diagonalization of the mass matrix $\mathcal{M}$ with the two states $\widetilde{\nu}_{R_1}$ and $\widetilde{\nu}^\dagger_{R_1}$ leads to eigenstates $\widetilde{N}_+$ and $\widetilde{N}_-$ with masses,

$$M_\pm \simeq M_1(1 \pm \frac{|B|}{2M_1}) \ , \tag{54}$$

where the leading order term $M_1$ is the F-term contribution from the superpotential (RH neutrino mass term) and the mass difference between the two mass eigenstates $\widetilde{N}_+$ and $\widetilde{N}_-$ is induced by the SUSY breaking $B$ term. The time evolution of the $\widetilde{\nu}_{R_1}$-$\widetilde{\nu}^\dagger_{R_1}$ system is governed by the Schrödinger

equation,

$$\frac{d}{dt}\begin{pmatrix} \widetilde{\nu}_{R_1} \\ \widetilde{\nu}^{\dagger}_{R_1} \end{pmatrix} = \mathcal{H}\begin{pmatrix} \widetilde{\nu}_{R_1} \\ \widetilde{\nu}^{\dagger}_{R_1} \end{pmatrix} , \tag{55}$$

where the Hamiltonian $\mathcal{H}$ is given by [26, 27],

$$\mathcal{H} = \mathcal{M} - \frac{i}{2}\mathscr{A} \tag{56}$$

$$\mathcal{M} = \begin{pmatrix} 1 & \frac{B^*}{2M_1} \\ \frac{B}{2M_1} & 1 \end{pmatrix} M_1 , \tag{57}$$

$$\mathscr{A} = \begin{pmatrix} 1 & \frac{A^*}{M_1} \\ \frac{A}{M_1} & 1 \end{pmatrix} \Gamma_1 . \tag{58}$$

For the decay of the lightest RH sneutrino, $\widetilde{\nu}_{R_1}$, the total decay width $\Gamma_1$ is given by, in the basis defined in Eq. (37) where both the charged lepton mass matrix and the RH neutrino mass matrix are diagonal,

$$\Gamma_1 = \frac{1}{4\pi}(\mathcal{Y}_\nu \mathcal{Y}^{\dagger}_\nu)_{11} M_1 = 0.374\ GeV . \tag{59}$$

The eigenstates of the Hamiltonian $\mathcal{H}$ are $\widetilde{N}'_{\pm} = p\widetilde{N} \pm q\widetilde{N}^{\dagger}$, where $|p|^2 + |q|^2 = 1$. The ratio $q/p$ is given in terms of $\mathcal{M}$ and $\Gamma$ as,

$$\begin{aligned} \left(\frac{q}{p}\right)^2 &= \frac{2\mathcal{M}^*_{12} - i\mathscr{A}^*_{12}}{2\mathcal{M}_{12} - i\mathscr{A}_{12}} \\ &\simeq 1 + Im\left(\frac{2\Gamma_1 A}{BM_1}\right) , \end{aligned} \tag{60}$$

in the limit $\mathscr{A}_{12} \ll \mathcal{M}_{12}$. Similar to the $K^0 - \overline{K}^0$ system, the source of CP violation in the lepton number asymmetry considered here is due to the CP violation in the mixing which occurs when the two neutral mass eigenstates ($\widetilde{N}_+$, $\widetilde{N}_-$), are different from the interaction eigenstates, ($\widetilde{N}'_+$, $\widetilde{N}'_-$). Therefore CP violation in mixing is present as long as the quantity $|q/p| \neq 1$, which requires

$$Im\left(\frac{A\Gamma_1}{M_1 B}\right) \neq 0 . \tag{61}$$

For this to occur, SUSY breaking, *i.e.* non-vanishing $A$ *and* $B$, is required. As the relative phase between the parameters $A$ and $B$ can be rotated away by an $U(1)_R$-rotaion, without loss of generality we assume from now on that the physical phase that remains is solely coming from the tri-linear coupling $A$.

The total lepton number asymmetry integrated over time, $\epsilon$, is defined as the ratio of difference to the sum of the decay widths $\Gamma$ for $\widetilde{\nu}_{R_1}$ and $\widetilde{\nu}^{\dagger}_{R_1}$ into final states of the slepton doublet $\widetilde{L}$ and the

Higgs doublet $H$, or the lepton doublet $L$ and the higgsino $\widetilde{H}$ or their conjugates,

$$\epsilon = \frac{\sum_f \int_0^\infty [\Gamma(\widetilde{\nu}_{R_1}, \widetilde{\nu}_{R_1}^\dagger \to f) - \Gamma(\widetilde{\nu}_{R_1}, \widetilde{\nu}_{R_1}^\dagger \to \overline{f})]}{\sum_f \int_0^\infty [\Gamma(\widetilde{\nu}_{R_1}, \widetilde{\nu}_{R_1}^\dagger \to f) + \Gamma(\widetilde{\nu}_{R_1}, \widetilde{\nu}_{R_1}^\dagger \to \overline{f})]} \tag{62}$$

where final states $f = (\widetilde{L}\ H)$, $(L\ \widetilde{H})$ have lepton number $+1$, and $\overline{f}$ denotes their conjugate, $(\widetilde{L}^\dagger\ H^\dagger)$, $(\overline{L}\ \overline{\widetilde{H}})$, which have lepton number $-1$. After carrying out the time integration, the total CP asymmetry is [26, 27],

$$\epsilon = \left(\frac{4\Gamma_1 B}{\Gamma_1^2 + 4B^2}\right) \frac{Im(A)}{M_1} \delta_{B-F} \tag{63}$$

where the additional factor $\delta_{B-F}$ takes into account the thermal effects due to the difference between the occupation numbers of bosons and fermions [32]. The final result for the baryon asymmetry is [26, 27],

$$\begin{aligned} \frac{n_B}{s} &\simeq -c\, d_{\widetilde{\nu}_R}\, \epsilon\, \kappa \\ &\simeq -1.48 \times 10^{-3} \epsilon\, \kappa \\ &\simeq -(1.48 \times 10^{-3}) \left(\frac{Im(A)}{M_1}\right) R\, \delta_{B-F}\, \kappa \end{aligned} \tag{64}$$

where $d_{\widetilde{N}}$ in the first line is the density of the lightest sneutrino in equilibrium in units of entropy density, and is given by, $d_{\widetilde{\nu}_R} = 45\zeta(3)/(\pi^4 g_*)$; the factor $c = (8N_F + 4N_H)/(22N_F + 13N_H)$ characterizes the amount of $B - L$ asymmetry being converted into the baryon asymmetry $Y_B$, with $N_F$ and $N_H$ being the number of families and the $SU(2)$ Higgs doublets, respectively. For the MSSM particle spectrum, $(N_F, N_H) = (3, 2)$. The parameter $\kappa$ is the dilution factor which characterizes the wash-out effects due to the inverse decays and lepton number violating scattering processes together with the time evolution of the system. It is obtained by solving the Boltzmann equations for the system. An approximation is given by [33]

$$10^6 \le r: \qquad \kappa = (0.1r)^{1/2} e^{-(\frac{4}{3})(0.1r)^{1/4}} \tag{65}$$

$$10 \le r \le 10^6: \qquad \kappa = 0.3/(r(\ln r)^{0.6}) \tag{66}$$

$$0 \le r \le 10: \qquad \kappa = 1/(2\sqrt{r^2+9})\ . \tag{67}$$

where $r$ is defined as

$$r \equiv \frac{M_{pl}}{(1.7)(32\pi)\sqrt{g_*}} \frac{(\mathcal{Y}_\nu \mathcal{Y}_\nu^\dagger)_{11}}{M_1} \tag{68}$$

with $M_{Pl}$ being the Planck scale taken to be $1.2 \times 10^{19}\ GeV$. We have $r = 183$ and correspondingly $\kappa = 0.00061$ in our model. The parameter $R$ is defined as the ratio,

$$R \equiv \frac{4\Gamma_1 B}{\Gamma_1^2 + 4B^2}\ , \tag{69}$$

which gives a value equal to one when the resonance condition, $\Gamma_1 = 2|B|$, is satisfied, leading to maximal CP asymmetry. As $\Gamma_1$ is of the order of $\mathcal{O}(0.1-1)\ GeV$, to satisfy the resonance condition, a small value for $B \ll \widetilde{m}$ is thus needed. Such a small value of $B$ can be generated by some dynamical relaxation mechanisms [34] in which $B$ vanishes in the leading order. A small value of $B \sim \widetilde{m}^2/M_1$ is then generated by an operator $\int d^4\theta Z Z^\dagger N_1^2/M_{pl}^2$ in the Kähler potential, where $Z$ is the SUSY breaking spurion field, $Z = \theta^2\ \widetilde{m} M_{pl}$ [27]. In our model, with the parameter $B' \equiv \sqrt{BM_1}$ having the size of the natural SUSY breaking scale $\sqrt{\widetilde{m}^2} \sim \mathcal{O}(1\ TeV)$, a small value for $B$ required by the resonance condition $B \sim \Gamma_1 \sim \mathcal{O}(0.1\ GeV)$ can thus be obtained.

Fig. 7 shows the ratio $R$ as a function of $B'$. For the specific value of the decay width $\Gamma_1$ predicted in our model, the resonance occurs at around $B' \sim 1.4\ TeV$. In Fig. 8, the region on the $Im(A)$ versus $B'$ plane that gives rise to an amount of baryon asymmetry consistent with the value derived from observation, $n_B/s = (0.87 \pm 0.04) \times 10^{-10}$, is shown. The required value for $B'$ near the resonance is around $800\ GeV - 2\ TeV$, and the required value for $|Im(A)|$ is around $(1-2)\ TeV$. At the resonance $B'$, the value for $|Im(A)|$ can be as low as $1\ TeV$ to generate sufficient amount of baryon asymmetry. In Fig. 9, we show the predictions for the asymmetry, $n_B/s$, as a function of $B'$ for different values of $Im(A)$. In the numerical analyses presented in Fig. 8 and 9, we assume $\delta_{B-F} = 1$. We note that even if an additional suppresion $\delta_{B-F} \sim 0.1$ is present, with a value of $Im(A) \simeq 10\ TeV$ at the resonance our model can still account for the observed BAU.

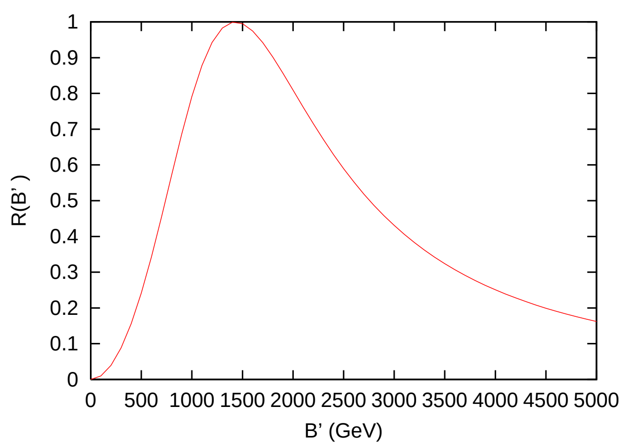


FIG. 7: The ratio $R$ as a function of $B'$. The resonance occurs at around $B' \sim 1.4\,TeV$.

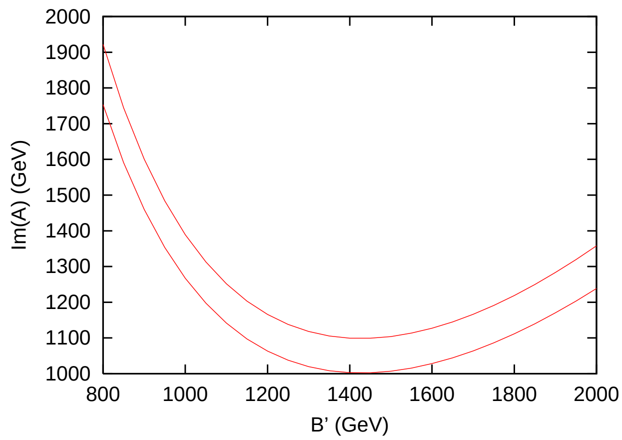


FIG. 8: The parameter space on the $Im(A)$ versus $B'$ plane that gives rise to an amount of baryon asymmetry consistent with the value derived from observations, $n_B/s = (0.87 \pm 0.04) \times 10^{-10}$, is the region bounded by these two curves. The upper curve corresponds to the upper bound from observation, $n_B/s = 0.91 \times 10^{-10}$, while the lower curve corresponds to the lower bound, $n_B/s = 0.83 \times 10^{-10}$.

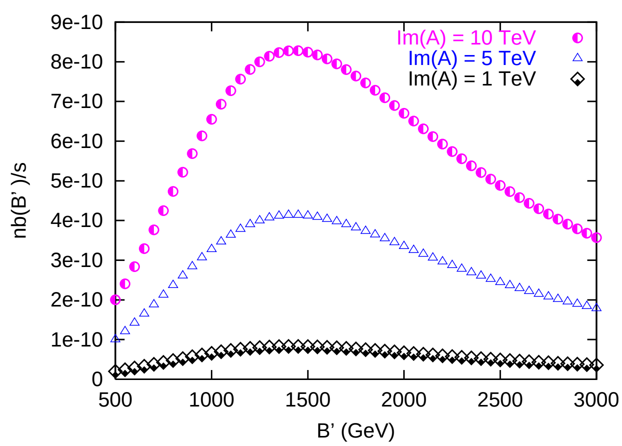


FIG. 9: The prediction for $n_B/s$ as a function of $B'$ for $|Im(A)| = 10\ TeV$ (circles), $5\ TeV$ (triangles) and $1\ TeV$ (squares).

## V. CONCLUSION

We have shown in this paper that, in contrast to the predictions of models with lop-sided textures, the predictions for LFV decays are well below the current experimental bounds. This is demonstrated in a model based on SUSY SO(10) with symmetric mass textures which give rise to predictions for all fermion masses and mixing angles, including those in the neutrino sector, that are in good agreement with experimental data within $1\sigma$. The predictions of our model for LFV processes, $\ell_i \to \ell_j \gamma$, $\mu - e$ conversion as well as $\mu \to 3e$, are well below the most stringent bounds up-to date. Our predictions for many processes are within the reach of the next generation of

LFV searches. This is especially true for $\mu - e$ conversion and $\mu \to e\gamma$. We have also investigated the possibility of baryogenesis resulting from soft leptogenesis. Our model predicts $M_1 < 10^9 GeV$ which is the required condition for this mechanism to work. With the soft SUSY masses assuming their natural values, $B' \sim 1.4\ TeV$ and $Im(A) \sim 1\ TeV$, we find that our model can indeed accommodate the observed baryon asymmetry of the Universe.

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